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LIGHT-PARTICLE EMISSION AT HIGH ANGULAR MOMENTUM

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1. Introduction

Emission of light particles from nuclear species at high angular momentum and excitation is a topic of long standing; yet, the discovery of deep-inelastic collisions has given a reviving impact to it as a possible probe of nuclear equilibration phenomena. Dissipation seems to start as soon as the two ions collide and may or may not result in a compound nucleus in which all degrees of freedom, including the shape, are in a state of thermal equilibrium.

It might be worthwhile to recall that our knowledge on nuclear behaviour at high angular momentum has almost entirely been collected by spectroscopy of cool nuclei. The highest levels from which transitions are resolved in γ spectra are within a very few MeV of the yrast line. Continuum γ rays go up to and somewhat beyond the particle binding above yrast. So, the interesting physics as far as experimentally accessible dwells in a narrow region above the yrast line where the lifetimes are long and measure by picoseconds. By all standards, such nuclei are in a state of perfect equilibrium with respect to energy, angular momentum of their constituent nucleons, and also with respect to their shape.

In contrast, particles are being emitted from regions tens of MeV above yrast where nuclear behaviour is described without gross inaccuracies by models which are either classical (Liquid Drop Model), or simple (Fermi Gas Model) - which merely states our yet insufficient knowledge. Trying to gain new experience, particle studies are done now at many laboratories, investigating the decay modes of nuclear systems with excitation energy above,100 MeV and at angular momenta approaching the limit of their stability.

Emission times for the first steps in the cascade are short enough to encounter nuclear systems remote from equilibrium. To give an example: at sufficiently high excitation, a damped binary system presumably undergoes particle decay faster than it takes to equilibrate its shape, so that there is the intriguing possibility to study the nascent, superdeformed stage of what will develop into a compound nucleus eventually.

The various stages of particle emission accompanying the evolution of the heavy-ion collision are sketched in Fig.1. Facing a highly complex situation, it is an all important question to which extent the experiment can provide us unique signatures. Such is probably the case for some of the very early stages, and as I shall show below, for the late stages: break up of the projectile at the very beginning, and decay from equilibrated fragments at the very end, both established by fragment-particle correlation experiments. In between, the formation of Volcov’s di-nuclear system (DNS) due to friction, as a quasi-stationary point.
in the evolution, might provide us with a meaningful distinction into
(i) emission during the approach phase til, say, the DNS has been formed with consider-
able equilibration regarding the N/Z ratio, the energy, the angular momentum, e.g., and
(ii) emission accompanying either the further evolution of the DNS into a compound
nucleus (shape equilibration), or accompanying its fragmentation into deep-inelastic
fragments and evaporation thereafter.

The first category comprises break up of the projectile with varying amount of partici-
aption of the target, including "deep-inelastic break up" and "pre-emission" as well as "promptly emitted particles" (PEP), "Fermi jets" as a physical consequence of
two interpenetrating Fermi gas containers, or potential wells. Experimental evidence
of these mechanisms is still in debate. The same is true for emission from "hot spots" or
hot zones - presumably on the borderline of categories (i) and (ii) - despite the
fact that this concept has found a remarkable echo in a number of studies. Category
(ii) comprises the recently proposed "pre-equilibrium evaporation" of the DNS, scis-
sion a's and evaporation from fragments.

In passing through this vocabulary, it appears that the role of high angular momen-
tum being of collective nature, is restricted to the later stages. However, high angular
momentum and pre-equilibrium emission might not be antagonistic concepts altogeth-
er, and following a suggestion of the organizers of this Conference, I shall try to
elaborate somewhat on this interesting perspective.

To be specific, I have chosen three chapters for my talk. The first deals with
particle evaporation from deep-inelastic fragments, to determine their spin and spin
alignment, and it is based on measurements a good part of which were done here in Stras-
bourg in collaboration with Dr. Scheibling and his group. A more sketchy discussion of
pre-equilibrium emission follows. In view of the strong demand for experimental time scales, the third part reports on an experiment designed to measure emission times for a highly excited compound nucleus, using particle-particle correlations.

2. Particle Evaporation from Deep-Inelastic Fragments: Spin and Alignment

I shall restrict the presentation of data to two systems: the light system 96 MeV
$^{14}$O+$^{58}$Ni studied at Strasbourg and Heidelberg and thoroughly investigated also
by circular polarization and fragment-neutron coincidence experiments; and the
system 400 MeV $^{40}$Ar+$^{93}$Nb studied at GSI.

A measurement of the intrinsic angular momenta of the fragments requires that
(i) the source is unique and identified, (ii) the source is in thermal equilibrium.
The sources are identified by the velocities of the particles detected at various detec-
tion angles, for fixed deep-inelastic binary Q value, mass or Z and angle of the light
fragment detected in coincidence. An example

![Figure 2: Velocity diagram for the 400 MeV $^{40}$Ar+$^{93}$Nb reaction, Q=-160 MeV, together with velocity spectrum of a particles detected at -48°](image-url)
Fig. 3. Differential α multiplicity $dM/dΩ$, the most-probable energy $E_α$, and a shape parameter $τ$ describing the fall-off in the α spectra, in the rest frame of the heavy fragment.

At backward angles opposite to the light-fragment detector (Ar type), α particles originate entirely from the heavy fragment (Nb type), and constancy of yield and shape of the spectra for $θ > 50^\circ$ ($80^\circ$ lab) indicates equilibration. The energy well agrees with Coulomb repulsion for Nb-fragment decay (the corresponding velocity is that of the upper circle in Fig. 2) and the shape parameter $τ$ is also consistent with this assumption (Fig. 3).

When describing the decay of equilibrated fragments in analogy to CN evaporation, an important difference is evident. It arises from the fact that DIC are mainly binary reactions and conservation laws do not-as in fusion reactions- play a decisive role. Rather, the sharing of energy, angular momentum and its substate population are the result of the very nature of the dissipative collision process implying fluctuations around average values.

The average angular momentum of the fragments can be classically considered to be induced by tangential friction forces with the tendency to reach rigid rotation, a concept which is also quantum mechanically meaningful as it marks the state of maximum level density $^1$ for frozen shape of the DNS. This "sticking limit" of the fragment spins is given by the moments of inertia $θ$ and the initial (and usually total) angular momentum $K_i$ as

$$J_{st} = K_i \theta / (θ_i + θ_f + μr^2)$$

for fragment no. 1, and $μr^2$ denotes the relative moment of inertia of the two fragments in the di-nuclear configuration.

The particle decay from a fragment with its spin $J_o$ pointing in the direction normal to the reaction plane will display an anisotropy in yield with respect to this axis. The physical reason of this anisotropy is the tendency towards maximum level density in the residual nucleus. Because of the exponential decrease of the density with spin, stretched transitions with $J_{st} \parallel J_o$ will be the most probable ones, and peaking therefore will occur in a plane orthogonal to $J_o$, in the reaction plane (Fig. 4).

The problem has been treated in the spirit of the statistical model in a semi-classical manner by Ericson and Strutinski back in 1958$^1$. A result more useful for coincidence studies where $J_o$ points in a fixed direction has been given by Th. Dössing$^1$:

$$W(θ) = \exp\left(\frac{-\hbar^2 (J_o + 1/2)^2}{2\Theta \cdot μr^2 (θ_i + θ_f + μr^2)} \sin^2 θ\right)$$

(1)

describes the angular distribution of the particles at angle $θ$ with respect to the spin axis of $J_o$, and the temperature $T$ and moment of inertia $θ$ refer to the residual nucleus. The expression is recognized as a Boltzmann factor in the rotational energy of the particle at geographical latitude $θ$ on the surface of the spinning nucleus$^{18}$. For practical purposes it should be noted that parameters $θ, T$ in a formally identical expression for a rotating Maxwell gas refer instead to the decaying nucleus.

Fig. 4 Geometry of particle decay. Spin $J_o$ perpendicular to reaction plane.
Let us see then the results of fragment-light-particle correlation experiments, in which the particle telescopes are moved out of the reaction plane (Fig. 5). For the 400 MeV Ar+ Nb reaction, the coincident out-of-plane yield for α's and protons is displayed in Fig. 6. θ_{α,p} = 90° denotes the reaction plane. Proton distributions are considerably flatter than those of α particles, which is expected from eq.(1), but in addition, protons come from all stages of the cascade, while α's are dominantly first chance. As there is no significant dependence on the in-plane angle, the data is summed, Fig. 7, and displays more clearly the small proton anisotropy.

Results for the 160+ 58Ni reaction, at 96 MeV, are shown in Fig. 8, protons and α's were detected in coincidence with C,N,O fragments at deep-inelastic Q values.
Our considerations of the particle decay up to now were based on a fixed spin axis. If the spins are distributed in direction, as a result of spin fluctuations, application of eq.(1) to data will yield values of $J_0$ too small. Let's first see the physics of these fluctuations and then look for an improved version of eq.(1).

Dissipation, we are taught by statistical mechanics, necessarily introduces fluctuations in the quantity considered. Nörenberg and Wolschin have applied such a transport description of deep-inelastic collisions also to angular momentum. Microscopically, fluctuations seem to arise, at energies below 20 MeV/u, not so much by nucleon-nucleon collisions, but rather by nucleons exchanged through the window of the contact zone between the interpenetrating ions which then collide with their additional momentum with the boundary of the partner. Other theories are reviewed in Ref.1. I mention this to underline that our ultimate aim in measuring macroscopic quantities like average fragment spin and alignment is to learn about the microscopic, or collective processes, by which colliding ions dissipate relative angular momentum. A more pedestrian view of spin fluctuations is depicted in Fig. 9.

![Fig.9. View of elementary angular momentum transfers, adding up in a more cooperative (a), or more random way (b).](image)

To account for this composite of correlated transfer and random walk in the parametrization of data, we choose a vector spin distribution consisting of a "macroscopic" spin $J_0$ aligned along the normal to the reaction plane and a fluctuating spin vector as a gaussian random variable with zero mean and variance $\sigma_J^2$. The average spin then is equal to $J_0$, and the resulting spin distribution is

$$P(\mathbf{J}) = \exp\left(-\frac{\langle \mathbf{J} - \mathbf{J}_0 \rangle^2}{2\sigma_J^2}\right).$$  \hspace{1cm} (2)

Folding $W(\theta)$ of eq.(1) with this spin vector distribution results in an expression for the out-of-plane distribution $W(\theta;\sigma_J^2)$ in which the exponent in eq.1 is modified by a factor

$$\left(1 + h^2\sigma_J^2/\Theta \cdot uR^2/(\mu R^2 + \Theta)^{-1}\right).$$ \hspace{1cm} (3)

A finite $\sigma_J$, or incomplete alignment, will decrease the anisotropy $W(90^\circ)/W(0^\circ)$, as expected intuitively.

We are now in need of an additional piece of information since $W(\theta)$ is not sufficient to determine both $J_0$ and $\sigma_J$. A somewhat lucky escape from this dilemma, for the data of Fig.7, is demonstrated in Fig. 10. Invoking that $J_0$ cannot exceed the sticking limit, we derive a meaningful upper limit $\sigma_J \leq 12 h$. Conversely, a lower limit of $J_0^2 \geq 28 h$ corresponds to complete alignment, $\sigma_J = 0$. The proton data are consistent regarding $J_0$, but much less conclusive regarding $\sigma_J$.

A more elaborate analysis was applied to the $0+\,\text{Ni}$ data, using first and second
moments of the $\gamma$ multiplicity\textsuperscript{12}. I have no time to enter the discussion of the somewhat tedious procedure to relate the multiplicity moments to those of the primary spin distribution\textsuperscript{13}. It is important, however, to realize that the width $\sigma |J|$ of the primary distribution of spin magnitudes, as derived from the variance of $M_{\gamma}$, relates to spin fluctuations only for light systems, like $O + Ni$ at 6 MeV/u where the $l_{+}$ window for deep-inelastic collisions is presumably very small. The multiplicity measurements are shown in Fig. 11.

Fig. 11. Average value and width of the $\gamma$ multiplicity for the system 96 MeV $^{16}O + ^{58}Ni$, in coincidence with carbon (○) and oxygen (○) fragments.

The combined analysis of particle and $\gamma$ data for $O + Ni$ uses a quantum mechanical expression for $W(\theta)$ devised by Dössing and described in Ref.13. The density matrix is assumed to be diagonal and given again by the spin distribution of eq.(2). The spin dependence of the $\alpha$ decay probability is taken from a statistical model calculation. The results are displayed in the $(J_0, \sigma_J)$ plane in Fig. 12. Three types of trajectories corresponding to the three experiments can be distinguished: the measured $a$ out-of-plane anisotropy, $A$, with its errors, defines a band swinging to the right and up for larger $\sigma_J$. The average of the spin magnitude $\langle |J| \rangle$, derived from $\langle M_{\gamma} \rangle$, decreases first slowly, then rapidly with $\sigma_J$. The experimental errors here are negligible, the uncertainties in deriving the primary spin considerable (hatched field). The widths of the spin magnitude distribution, from $\sigma_J$, yield trajectories running almost vertically. Only the region between the inner two lines is due to the experimental error.

The result is not very comforting as the cores of the three branches don't overlap. Although, I think, the uncertainties which come especially into the $\gamma$ data analysis have been generously assessed, the poor consistency might just point to them. A more interesting version is given below.

Fig. 12. Trajectories for combined analysis of particle and $\gamma$ data, see text. 96 MeV $^{16}O + ^{58}Ni$, a particles and $\gamma$ rays coincident with oxygen fragments.

An equivalent description of the spin distribution is provided by the alignment parameter\textsuperscript{16}

\[ P_{zz} = \frac{3}{2} <J_z^2> - \frac{1}{2} = \frac{3}{2} <\cos^2 \theta> - \frac{1}{2} \]

expressing the semi-classical tilt angle. The $P_{zz}$ values for the data presented range between 0.85 and 0.8.

It is interesting which equilibrium, or long-time limit, is to be expected for $\sigma_J$. Unfortunately, no expression for one particular fragment is given in the literature; those given only apply to the DNS before
separation into the fragments. A simple equipartition estimate, assuming 3 independent rotational degrees of freedom, yields $a = 30T$, which amounts to $a = 8.8\hbar$ and $17\hbar$ for Ni and Nb, respectively. A probably much better assumption is that the fragments in equilibrium form a rigid rotor which adjusts to a statistical K distribution. The spin distribution of the fragments is then

![Fig. 13. K distribution of the DNS in equilibrium (a) and resulting fragment spin distribution after scission along the symmetry axis (b).](image)

obtained by use of the standard assumption that the complex scissions along the symmetry axis. An elaborate discussion of the long-time limit of $Q$ involving various collective degrees of freedom, like bending and wriggling modes, has been given by Moretto and Schmitt. These modes result in additional spin fluctuations.

I have the strong impression that an important dynamical aspect has been neglected so far. Within the rotating DNS, the fragments are exposed to a large inertial or centrifugal force which tends to disrupt the system in a direction perpendicular to the rotation axis ($J$ or $z$ axis in Fig. 13). In consequence, the fragment spin distribution, after separation of the complex, will be re-aligned. Such an effect is of mere academic interest for low-energy fission, and not mentioned in the book; still small for sequential fission of deep-inelastic fragments, it grows to size for the large angular momenta of a DNS, as a simple consideration of the balance of forces (Fig. 14) shows: For the example of the 400 MeV Ar$^+$ Nb reaction, a DNS formed by $J = 120\hbar$ (=$0.8I_\hbar$) and assumed to carry an alignment of $P_{zz} = 0.6$ (=$\theta_{rms} = 30^\circ$) only, will by centrifugal re-alignment result in a highly aligned fragment spin distribution, with $P_{zz} = 0.9$ (=$\theta_{rms} = 16^\circ$). I note here in parenthesis that the combined experimental results analyzed in Fig. 12 might be taken as evidence that the $x,y$ components of $\sigma$ are significantly smaller than its $z$ component, due to this effect.

The experimental determination of the size of spin fluctuations and the question of their physical origin is an important matter. I have to conclude that it remains completely open today, whether measured variances, including those from sequential fission studies, correspond to equilibration in some way or the other, with regard to which collective degrees? Or whether variances are enlarged by non-equilibrium contributions and quantum fluctuations. Anyway, progress apparently needs both, improving our experimental skill and further theoretical guidance.

3. Fast Particle Emission: A Consequence of High Angular Momentum?

The idea is not as curious as it might appear, as we are aware that deep-inelastic collisions themselves proceed faster than compound-nucleus formation due to lack of stability for high partial waves (J. Wilczynski). Why should we not encounter similar physics in particle emission? I like to present some evidence that we actually do.

In order to avoid confusion, let me first remind you of the overall pre-equili-
brium pattern as it is observed in collisions with sufficiently fast projectiles, again for the example of 400 MeV $^{40}$Ar+$^{93}$Nb, Ref. 27, which is shown in Fig. 15 for a particles in coincidence with Ar-type fragments.

![Figure 15](image1.png)

**Fig. 15.** In-plane correlation of a particles with light fragments detected well behind the grazing angle at position labeled H.I. $\theta = 0$ refers to recoil direction of heavy fragment, angle and yield in this rest frame.

Events are deep-inelastic ($Q^e = -160$ MeV) and originate from Nb-type fragments, with some uncertainty close to the beam direction. The $\alpha$ channel seems to be the most favourable one to observe pre-equilibrium emission, as there is little multiplicity from later, equilibrated stages of the cascade.

This pre-equilibrium emission has many similarities to the one observed in the 96 MeV O+Ni system which was interpreted by Ho et al. 9 as evidence for emission from a hot zone. There is considerable interest to look for a "shadow" in the angular correlation, on the side of the light-fragment detector, as this is considered to be a strong evidence for a hot spot mechanism. 28. Possibly, such a shadow has been observed for the 101 MeV $^{16}$O+Ti system by the Strasbourg group 29, and I like to show their data in Fig. 16. The correlation show their data in Fig. 16. The correlation is for a selection of deep-inelastic events for which the $^{16}$O excitation remains below the effective $\alpha$ threshold. If substantiated, the shadow is indeed remarkably deep and sharp.

I like to convince you now, that we observe two different mechanisms in the Ar+Nb reaction. Evidence for a "direct" component evolves from close inspection of a velocity spectra taken at forward lab angles, where a shoulder emerges which has about beam velocity (Fig. 17). Gating on the "direct" $\alpha$'s does not change the energy distribution of the coincident projectile-like fragments in a significant way. This invites to interpret these $\alpha$'s as being emitted prior to deep-inelastic scattering of the remainder of the projectile. 27. Much further support to this conjecture is given by a-light fragment correlation experiments probing the 148 MeV $^{15}$N+$^{58}$Ni system, by the Birmingham group 30, which are presented at
Fig. 17. Coincident α velocity spectra at 4°, 37.5°, and 3° lab angles revealing a "direct" component of about beam velocity (full line). 400 MeV \(^{9}Ar^{+}\)\(^{59}Nb\), geometry as in Fig. 2. Coulomb escape velocity from Nb-like fragment is also shown (dashed line).

Fig. 18. In-plane correlation of α particles with various light fragments for the system 148 MeV \(^{11}Na^{+}\)\(^{58}Ni\). Solid curves are derived from singles α angular distributions. From Ref. 30.

this Conference, and part of the data is shown in Fig. 18. The authors show that the coincident cross section for α particles well factorizes into a product of the singles cross section for α emission and the singles heavy-ion cross section. They infer that the α particles are emitted at an early stage, prior to the formation of the deep-inelastic fragments.

Fig. 19 depicts the concept of what might be called the "incomplete collision": the emission of an α particle with momentum close to what it was inside the projectile precedes the scattering-or fusion-of the remainder of the projectile. While the α particle is more of a spectator, the projectile residue experiences various degrees of energy and angular momentum dissipation.

The incomplete fusion, observed already by Sikkeland et al.\(^3^1\), has recently come intensely into discussion, stimulated by very informative experiments\(^3^2\). At the early side, quasi-elastic break-up has also been studied recently by correlation experiments, as mentioned above.

Its direct character quite obviously established by the large forward momentum, it is interesting to regard the fast α emission from a phase space point of view. Regarding the DNS as a stationary point in the evolution, its available level density is the larger, the smaller the amount of energy is which is tied up in collective rotation.

Fast α emission, removing a large amount of angular momentum from the system, is therefore statistically favored. To show that it is also dynamically favored, we draw a
potential diagram (Fig. 20) for the somewhat simplistic configuration of an α particle riding on the outer tip of the rotating dinuclear system. It is apparent that the α

\[
V \quad (\text{MeV})
\]

\[
\alpha(\ell) + V_{\text{c}}(\ell)
\]

\[
V_{\alpha}(\ell) + V_{\text{c}}(\ell)
\]

\[
V_{\alpha}(\ell) + V_{\text{c}}(\ell)
\]

\[
\frac{R_\alpha(\lambda)}{R_\alpha(1)} = P_\alpha(\lambda) / P_\alpha(1)
\]

Fig. 20. Potential diagram for an α particle in configuration depicted for a grazing orbit. Coulomb and centrifugal potentials superimposed in the upper curve. Sketch of situation for 400 MeV \(^{40}\text{Ar} + ^{93}\text{Nb}\).

becomes unbound already by the Coulomb potential of the nearby target, but even more so due to the large centrifugal potential felt in its orbit around the common centre of mass. The centrifugal potential probably builds up very rapidly as soon as contact forces between the colliding partners become effective, much earlier than rigid rotation sets in. A rough estimate for this case, of the relative α escape probability as a function of the angular momentum brought into the collision, is given in Fig. 21.

Despite their simplicity, these arguments were given in order to show that "direct" emission patterns might also evolve from a statistical and dynamical consideration of the very specific initial configuration encountered in deep-inelastic collisions being characterized by very high angular momentum and highly non-spherical boundaries.

Maybe, a "constrained phase-space analysis" as very successfully applied\(^{33}\) to direct multinucleon transfer reactions, could provide a sound basis.

It is very interesting to follow this concept alongside the evolution of the DNS. Because of large viscosity, its shape adjusts only slowly (\(\tau_{\text{shape}} = 2.4 \times 10^{-2}\) s, Ref.1) — if at all — to the potential minimum of equilibrium shape of the compound nucleus, at given angular momentum.

In Fig. 22, I show a schematic diagram of two Yrast lines, one for the di-nuclear system, and one for the compound nucleus, supposedly in its equilibrium shape. The two

Fig. 21. Estimate of relative α escape probability, combining phase space (dashed) and barrier penetration, as function of \(\lambda\). Relates to Fig. 20.

Fig. 22. Yrast lines for DNS and CN, in anticipation of the situation for 400 MeV \(^{40}\text{Ar} + ^{93}\text{Nb}\). See text.
yrast lines are displaced by an amount $Q$ which is the energy set free - for the light system under consideration - when the two ions at contact are fused into a spherical, non-rotating compound-nucleus. The arrow in Fig.22 indicates, that particle emission again is favoured by phase space, and that it is expected to strongly enhance (incomplete) fusion of high partial waves.

Pohlhofer has recently discussed his evaporation residue data for the reaction

![Fig.23. Mass distributions of evaporation residues for the reaction $^{46,52,56}$Ti + $^{24}$Mg $\rightarrow ^{79}$Ge at three energies (histogram) compared to CASCADE calculations without deformation. (heavy lines). From Ref. 34.](image)

$^{46,52,56}$Ti + $^{24}$Mg in terms of "pre-equilibrium evaporation", i.e. for the rotating, highly deformed DNS. Such, he seems to explain the large yield for multiple $\alpha$ emission which cannot be accounted for by evaporation calculations without gross deformation.

Superdeformation of compound nuclei induced and together with high angular momentum, eventually, has been shown by Blann et al.\textsuperscript{25} to enhance $\alpha$ emission over fission, and the results of the calculations for $^{149}$Tb compound nuclei are shown in Fig.24. No better conditions for such mechanisms to become effective can be imagined than the DNS, in thermal or angular momentum equilibrium or not, but prior to shape equilibration.

![Fig.24. Partial decay probabilities for $\alpha$, $n$, $p$ and including fission (f), for $^{149}$Tb at 120 MeV of excitation, as function of I. Calculations for spherical (full) and deformed nuclei (broken). From Ref. 35.](image)

On occasion of this Conference, I tried to put the discussion of the high-angular momentum aspect in pre-equilibrium particle emission into the foreground; a ground not always quite firm is probably typical of a novel and exciting field.

Two-Particle Correlations: A Method to Measure Compound-Nucleus Emission Times

As the importance of time scales all along the heavy-ion collision is quite evident, I like to close my talk by presenting an experiment which is devised to measure emission times for a highly excited CN, in the range below $10^{-28}$s.

The experiment of Aichelin et al.\textsuperscript{36} was born out of the idea to exploit intensity correlations, or, the quantum mechanical interference of identical-particle amplitudes\textsuperscript{37}, by detecting two like
particles in coincidence from an incoherent source (CN). It turned out, however, that the final-state interaction (FSI) between, say, two α particles or two protons, is by far dominant, and the quantum interference was not observed.

FSI might not at all turn out to be inferior as a tool to measure lifetimes: it is the effective distance \( r_{12} \) between the particles in the neighbourhood of the source which decides upon whether the interaction between the particles \( V(r_{12}) \), Coulomb and nuclear, is effective in distorting particle spectra and yields. For a CN of lifetime \( \tau \), and assuming here that \( \tau \) is the same for the daughter nucleus, the mean separation is equal to the distance travelled by the first particle during \( \tau \)

\[ <r_{12}> = \tau <v>. \tag{4} \]

The experiment was done at the Heidelberg post-accelerator\(^3\) by choosing the reaction \( ^{14} \text{He} + ^{76} \text{As} \rightarrow ^{16} \text{O} + ^{72} \text{Ge} \), achieving 122 MeV of excitation in this CN.

Denoting the coincidence yields, or spectra, by \( C_{12} \) and \( C_{13} \), respectively, CN symmetry asserts that \( C(\phi) = C(\pi - \phi) \) in the centre of mass, and therefore \( C_{12} = C_{13} \) in the laboratory system. The last relation does not hold exactly since the recoil of the daughter nucleus enhances the solid angle for pair (1,3) by a small amount over that of pair (1,2) in the centre of mass. The detector assembly actually used in the experiment is shown in Fig. 26. Particles were identified by their time of flight (pulsed beam), to have precise energy calibration.

It is essential in such an experiment to supply simultaneously a measurement of two-particle spectra undistorted by the effect to be observed, under otherwise identical conditions. For CN decay, the geometry shown schematically in Fig. 25 comes close to this ideal. Detectors are arranged on a cone of constant scattering angle \( \theta \), assuring identical singles spectra; the "signal" expected to show up in the close pair (1,2) may be compared to the "background" in pair (1,3).

The two-particle spectra are displayed as a function of the difference energy \( \Delta E = E_{11} - E_{12} \), say, with no window on the sum energy. Results for the far geometry (1,3) are displayed in Fig. 27. The spectral shape is indicative of two independent evaporation...
totally random spectrum (right) reveals better than any calculation that a compound nucleus at such high excitation energy behaves in the first place like an oven. So, Fig. 27, left, provides the reference or background spectrum, and it originates apparently from a highly statistical source.

In Fig. 28 two difference-energy spectra measured in close geometries are shown in comparison. The dominant feature are two peaks symmetric to zero difference energy, which originate from break up of $^6$Be(g.s.). Due to the long-lived ground state and the small Q value of 92 keV, these peaks move as a function of $\Delta \phi$ in a very distinctive way, and disappear for $\Delta \phi > 30^\circ$.

![Fig. 28. aa' difference-energy spectra for two close geometries (top and middle), together with reference spectrum. Arrows label peak in \( C_{13}(176^\circ) \) for comparison of relative cross sections.](image)

![Fig. 29. Plot of the quantity \( S_{12}(\Delta E) \), defined in text, for aa and pp correlations in \( 4^\circ \) geometry. Offset merely to avoid neg. numbers.](image)

Before returning to the possible origin of the $^6$Be(g.s.) events, the smooth part of the spectra under the peaks deserves a closer inspection. The data of Fig. 28 show that the absolute yield $C_{12}(4^\circ)$ is only about half of $C_{13}(176^\circ)$ at $\Delta E = 0$, despite the tails from the Be peaks.

In order to do a close comparison independent from possible normalization errors and derive the "signal" from the "noise", we use the quantity

\[
S_{12}(\Delta \phi, \Delta E) = \ln \left( \frac{C_{12}(\Delta \phi, \Delta E)}{C_{13}(176^\circ, \Delta E)} \right)
\]

plotted in Fig. 29. The depletion of the spectrum measured in the close geometry is clearly visible, and the depletion seems to reach its maximum for $\Delta E = 0$ in the aa as well as in the pp correlation.
The observed effects are by far too large, and of wrong $\Delta E$ dependence, to be accounted for by recoil effects. The depletion of the small-angle $\alpha\alpha$ correlation, clearly, is just opposite to a constructive boson interference.

We are not aware of any other example of a statistical nuclear reaction in which effects of the mutual interaction between two emerging particles have been observed. On the other hand, the environment of a CN emitting, even at very high excitation energy, particles with small kinetic energies from its surface, favours the low relative momenta at which the attractive s-wave nuclear interaction dominates. Here, $^8\text{Be}\, (g.s.)$ comes into play, in a twofold way: first, G. Bertsch and the calculated formation probability for Be, given elsewhere, are consistent with the lifetimes estimated below. Here, I like to give only a line of thought, rather than precise results, how the lifetime can be derived from the measured depletion pattern, along with questions (i) and (ii).

(i) For coalescence to occur, the two $\alpha$ particles must enter a sphere of radius $R_{\text{Be}}$. For this to happen, the first-emitted particle is allowed to travel at most a distance $R_{\text{Be}}$ before the second is emitted. Denoting the average velocity at the surface by $<u>$, the lifetime is

$$\tau \approx \frac{R_{\text{Be}}}{<u>},$$

and since the average kinetic energy at the surface is equal to $T$, we obtain for $R_{\text{Be}} \approx 4\,\text{fm}$ and $T = 3.5\,\text{MeV}$, $	au \approx 3.1 \times 10^{-22}\,\text{s}$.

(ii) For coalescence to occur, the two $\alpha$ particles must enter a sphere of radius $R_{\text{Be}}$. For this to happen, the first-emitted particle is allowed to travel at most a distance $R_{\text{Be}}$ before the second is emitted. Denoting the average velocity at the surface by $<u>$, the lifetime is

$$\tau \approx \frac{R_{\text{Be}}}{<u>},$$

and since the average kinetic energy at the surface is equal to $T$, we obtain for $R_{\text{Be}} \approx 4\,\text{fm}$ and $T = 3.5\,\text{MeV}$, $	au \approx 3.1 \times 10^{-22}\,\text{s}$.

Before trying to answer (ii), let us look at the $^8\text{Be}\, (g.s.)$ wave function in Fig. 31. It is large inside $R_{\text{Be}}$ and very small outside, in response to the extremely narrow resonance width of 8.8 eV (Fig. 31). The strong rise of $\phi(r)$ towards the interior results in a momentum distribution $\Delta p \approx \frac{\pi}{R_{\text{Be}}} \approx 50\,\text{MeV/c}$ wide. In order to synthesize, the two $\alpha$'s must not exceed 50 MeV/c relative momentum. A relative longitudinal momentum of this size corresponds to a most probable kinetic energy difference, at the surface, of 300 KeV only - one $\alpha$ being actually at rest, the other going radially with 300 keV. As kinetic energy differences are maintained during acceleration in the Coulomb field of the nucleus, it also follows that $\Delta E \lesssim 300\,\text{keV}$ (neglecting $\alpha\alpha$ Coulomb repulsion!) is the condition for Be formation in experiment. This, however, is much less
less than the range in $\Delta E$ over which depletion of the smooth $aa$ background is observed (Fig. 29, top).

So far, in discussing an answer to (ii), we have neglected the finite lifetime of the CN. Actually, when we observe two particles with energies $E_A$ and $E_B$ in the Si detectors, their energies at the nuclear surface are not known exactly, but described by a random probability distribution, presumably a Lorentzian of width $\Gamma = \hbar/\tau$. Such, as the short-lived nucleus emits wave packets of width $\Gamma$, formation of Be is enhanced, because the interval of difference energies which are acceptable to the momentum condition spreads out to the order of $\Gamma$ (Fig. 32). Correspondingly, the depletion pattern widens.

![Fig. 32.](image-url) Particles detected with energies $E_A$, $E_B$ have identical energies at the nuclear surface to the extent their wave packets overlap, which depends on the lifetime $\tau = \hbar/\Gamma$ of the CN.

Let us interpret the $aa$ data in this way and obtain a second estimate of the lifetime. The two probability distributions of Fig. 32 have to be folded with the momentum distribution in $^9$Be(g.s.) - the Fourier transform of $\phi(r)$ in Fig. 31 - and averaged over and weighted by the evaporation spectra at the nucleus. Then, one obtains a window $W$ of difference energies allowed by the momentum condition $\Delta p_{rel} < 50$ MeV/c, as a function of $\Gamma$. Although straightforward, it cannot be done in a closed expression even for gaussians. A reasonable estimate probably is $W = 3\Gamma$. With $W = 8$ MeV, from Fig. 29 we obtain

$$\tau = \frac{\hbar}{\Gamma} = 3\hbar/\Gamma W = 2.5 \times 10^{-22}\text{s}.$$

The lifetime estimates should be taken with some caution as the inter-particle Coulomb effects, especially, have been neglected. However complicated a full analysis of such measurements eventually will turn out, it will be essentially model-independent and (only) be the necessary and highly desired means to translate measured correlation spectra into lifetimes. Actually, even the measurement of decay curves can be envisioned, as the depletion spectrum, and its dependence on $\Delta \phi$, contains much more information, than just the mean lifetime. I am optimistic that this aim will be achieved, but until then, the question mark in the heading to this section, better remains.

5. Conclusion

It is probably fair to state that particle emission is a tool well chosen to study angular momentum effects as well as characteristic time scales in the evolution of heavy-ion collisions. Some of it is promise, other parts have already come under detailed investigation. In both respects, a lot remains to do. For topics to which I had either no time or no experience to turn to, I like to refer the reader of this volume to the talks of Dr. D. Guerreau and Dr. F. Plasil.

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